STUDY THE EXOTIC Θ^+ IN POLARIZED PHOTOPRODUCTION REACTIONS

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We present an analysis of a Beam-Target double polarization asymmetry in $\gamma n \to \Theta^+ K^-$. We show that this quantity can serve as a filter for the determination of the Θ^+'s spin-parity assignment near threshold. It is highly selective between $1/2^+$ and $1/2^-$ configurations due to dynamical reasons.

1. Introduction

The report of signals of a strangeness $S = +1$ baryon Θ^+ [1] has stimulated a tremendous number of activities in both experiment and theory in the past 12 months. Baryon with strangeness $S = +1$ and strongly coupled into $\bar{K}N$ implies that its minimum number of constituents should be at least five, i.e. a pentaquark state ($uudd\bar{s}$). With a rather-low mass of 1.54 GeV and a narrow width of less than 25 MeV, the existence of such a state (later another state Ξ^{--} was also reported [2]) also initiates a tremedous number of questions about strong QCD dynamics in the low energy regime. Whether it is the chiral-soliton-model-predicted $\bar{\mathbf{10}}$ multiplets [3], or pentaquark $\bar{\mathbf{10}}$ in the quark model [4,5], such a state reveals some important new aspects of QCD, which have not been widely concerned before.

The quantum numbers of the Θ^+ are undoubtedly one of the key issues for answering some of those critical questions. So far, the spin and parity of the Θ^+ have not been well-determined in experiment [6,7], in spite of that the absence of signals in $\gamma p \to \Theta^{++} K^-$ favors Θ^+ to be an isopsin singlet (see e.g. SAPHIR results). Meanwhile, a series of theoretical approaches have been proposed to understand the nature of the Θ^+. Due to the limited space here, we will concentrate on the theoretical study of the Θ^+ photoproduction on the nucleon. A number of phenomenological approaches have been proposed in the literature focussing on cross section

predictions [8,9]. However, due to the lack of knowledge about the underlying reaction mechanism, such studies of the reaction cross sections are strongly model-dependent. For instance, the total width of Θ^+ still has large uncertainties and could be much narrower [10,11,12], and the role played by K^* exchange, as well as other s- and u-channel processes are unknown. Also, in a phenomenological approach the energy dependence of the couplings is generally introduced into the model via empirical form factors, which will bring further uncertainties. Taking this into account, there are advantages with polarization observables (e.g. ambiguities arising from the unknown form factors can be partially avoided). In association with the cross section studies, supplementary information about the Θ^+ can be obtained [13,14,15,16,17].

2. A minimum phenomenology

The following effective Lagrangians are assumed for the ΘNK couplings with the spin-parity $1/2^+$ and $1/2^-$ for the Θ^+, respectively:

$$\begin{aligned} L_{eff}(1/2^+) &= g_{\Theta NK}\bar{\Theta}\gamma_\mu\gamma_5\partial^\mu KN + \text{h.c.}, \\ L_{eff}(1/2^-) &= g_{\Theta NK}\bar{\Theta}NK + \text{h.c.}, \end{aligned} \tag{1}$$

where $\bar{\Theta}$, N and K denote the field of Θ^+, neutron and K^-. The coupling constant $g_{\Theta NK}$ is determined by the experimental data for the decay width of $\Theta^+ \to NK$, which however is still very imprecise. In this work, the same set of parameters as Refs. [13,14] is adopted by assuming $\Gamma_{\Theta^+\to NK} \simeq 10$ MeV. In Ref. [18], a narrower width of 1 MeV is adopted, based on analyses of N-K scattering data [10]. The change of such an overall factor will not change the behavior of polarization asymmetries for exclusive Born terms. The magnetic moment of Θ^+ is estimated based on the phenomenology of diquark model of Ref. [4], which is consistent with other models in the literature [19,9].

We also include the K^* exchange in this model as the leading contribution in association with the Born terms. The $K^*N\Theta$ interaction is given by

$$L_{\Theta NK^*}(1/2^+) = g_{\Theta NK^*}\bar{\Theta}(\gamma_\mu + \frac{\kappa_\theta^*}{2M_\Theta}\sigma_{\mu\nu}\partial^\nu)V^\mu N + \text{h.c.} \ , \tag{2}$$

and

$$L_{\Theta NK^*}(1/2^-) = g_{\Theta NK^*}\bar{\Theta}\gamma_5(\gamma_\mu + \frac{\kappa_\theta^*}{2M_\Theta}\sigma_{\mu\nu}\partial^\nu)V^\mu N + \text{h.c.} \ , \tag{3}$$

where $g_{\Theta NK^*}$ and κ^*_θ denote the vector and "anomalous moment" couplings, respectively. We follow Refs. [13,14] to adopt the values for $g_{\Theta NK^*}$ and κ^*_θ. It should be note that $|g_{\Theta NK^*}| = |g_{\Theta NK}|$ and $\kappa^*_\theta = 0$ are widely adopted in the literature. In Ref. [18], $g^2_{\Theta NK^*}(1/2^+) = 3g^2_{\Theta NK}$ and $\kappa^*_\theta = 0$ are adopted in the calculation of cross sections for $1/2^+$ pentaquarks with the consideration of the "fall-apart" mechanism [20,12,21,22]. The change of parameters for the K^* exchange does not change the behavior of the BT asymmetries for $\gamma n \to \Theta^+ K^-$ near threshold due to the importance of the Born terms.

The effective Lagrangian for the $K^*K\gamma$ vertex is given by

$$L_{K^*K\gamma} = \frac{ie_0 g_{K^*K\gamma}}{M_K}\epsilon_{\alpha\beta\gamma\delta}\partial^\alpha A^\beta \partial^\gamma V^\delta K + \text{h.c.} \, , \tag{4}$$

where V^δ denotes the K^* field. The coupling $g_{K^{*0}K^0\gamma} = 1.13$ is determined by $\Gamma_{K^{*0}\to K^0\gamma} = 117$ keV [23].

In the photoproduction reaction, defining the z-axis as the photon momentum direction, and the reaction plane in x-z in the c.m. system of γ-n, the transition amplitude for $\gamma n \to K^-\Theta^+$ can be expressed as

$$T_{\lambda_\theta,\lambda_\gamma\lambda_N} \equiv \langle \Theta^+, \lambda_\theta, \mathbf{P}_\theta; K^-, \lambda_0, \mathbf{q}|\hat{T}|n, \lambda_N, \mathbf{P}_i; \gamma, \lambda_\gamma, \mathbf{k}\rangle, \tag{5}$$

where $\lambda_\gamma = \pm 1$, $\lambda_N \pm 1/2$, $\lambda_0 = 0$, and λ_θ are helicities of photon, neutron, K^-, and Θ^+, respectively.

3. Double polarization asymmetries

We are interested in a Beam-Target (BT) double polarization asymmetry measured in $\gamma n \to \Theta^+ K^-$. In this reaction the photons are circularly polarized along the photon moment direction $\hat{z}$ and the neutrons transversely polarized along the $\hat{x}$-axis within the reaction plane. In terms of the density matrix elements for the Θ^+ decays, the BT asymmetry can be expressed as,

$$D_{xz} = \frac{\rho^{BT}_{\frac{1}{2},\frac{1}{2}}}{\rho^0_{\frac{1}{2},\frac{1}{2}}}, \tag{6}$$

where the subscript xz denotes the polarization direction of the the initial neutron target along x-axis in the production plane and the incident photon along the z-axis, and the definition of the density matrix element is

$$\rho^{BT}_{\lambda_\theta\lambda'_\theta} = \frac{1}{2N}\sum_{\lambda_\gamma,\lambda_N} \lambda_\gamma T_{\lambda_\theta,\lambda_\gamma-\lambda_N} T^*_{\lambda'_\theta,\lambda_\gamma\lambda_N}, \tag{7}$$

where, $N = \frac{1}{2}\sum_{\lambda_\theta,\lambda_\gamma,\lambda_N} |T_{\lambda_\theta,\lambda_\gamma\lambda_N}|^2$ is the normalization factor.

As follows, we express the transition amplitudes in terms of the CGLN amplitudes [24]. This is useful for our understanding the behavior of the BT asymmetry near threshold.

For $1/2^+$, we have

$$\langle \Theta^+, \lambda_\theta, \mathbf{P}_\theta; K^-, \lambda_0, \mathbf{q}|\hat{T}|n, \lambda_N, \mathbf{P}_i; \gamma, \lambda_\gamma, \mathbf{k}\rangle = \langle \lambda_\theta|\mathbf{J}\cdot\boldsymbol{\epsilon}_\gamma|\lambda_N\rangle \ , \tag{8}$$

with

$$\mathbf{J}\cdot\boldsymbol{\epsilon}_\gamma = if_1\boldsymbol{\sigma}\cdot\boldsymbol{\epsilon}_\gamma + f_2\frac{1}{|\mathbf{q}||\mathbf{k}|}\boldsymbol{\sigma}\cdot\mathbf{q}\boldsymbol{\sigma}\cdot(\mathbf{k}\times\boldsymbol{\epsilon}_\gamma) + if_3\frac{1}{|\mathbf{q}||\mathbf{k}|}\boldsymbol{\sigma}\cdot\mathbf{k}\mathbf{q}\cdot\boldsymbol{\epsilon}_\gamma + if_4\frac{1}{|\mathbf{q}|^2}\boldsymbol{\sigma}\cdot\mathbf{q}\mathbf{q}\cdot\boldsymbol{\epsilon}_\gamma. \tag{9}$$

The coefficients $f_{1,2,3,4}$ are functions of energies, momenta, and scattering angle $\theta_{c.m.}$, and contain information on dynamics. They provide an alternative expression for the BT asymmetry:

$$D_{xz} = \sin\theta_{c.m.}\mathrm{Re}\{f_1f_3^* - f_2f_4^* + \cos\theta_{c.m.}(f_1f_4^* - f_2f_3^*)\}. \tag{10}$$

We are interested in the energy region near threshold, namely, with the overall c.m. energy $W \sim 2.1$ GeV. This is the region that the well-established contact term (Kroll-Ruderman term) dominates over all the other processes, and is the main component of f_1. The term of f_2 will have contributions from the s- and u-channel, while the terms of f_3 and f_4 from the t- and u-channel. Near threshold, it shows that the t-channel amplitudes have the least suppressions from the propagator $1/(t - M_K^2)$ in comparison with the s- and u-channel. Also, the terms contributing to f_4 will be further suppressed near threshold due to the small momentum $|\mathbf{q}|$ in the final state. Considering the products of f coefficients in Eq. (10), we find the dominant contributions are from:

$$f_1f_3^* \simeq -e_0^2g_{\Theta NK}^2\frac{2}{t - M_K^2}F_c(k,q)F_t(k,q) \ , \tag{11}$$

where $F_c(k,q)$ and $F_t(k,q)$ are form factors for the contact and t-channel, and are treated the same in this approach.

This kinematic analysis leads to the robust prediction of the behaviour of D_{xz} near threshold to be dominated by

$$D_{xz} \simeq \sin\theta_{c.m.}\mathrm{Re}\{f_1f_3^*\}. \tag{12}$$

Since the CGLN coefficients only depend weakly on $\theta_{c.m.}$ (via the Mandelstam variables), the dominance of the above term implies a $\sin\theta_{c.m.}$ behavior of D_{xz}, and the sign of D_{xz} is determined by the product.

In Fig. 1, the numerical results for the BT asymmetry at $W = 2.1$ GeV are presented. The solid curve denote results in the Born limit, while the

dashed and dotted curves represent results for including the K^* exchange. Although the sign change of the K^* exchange results in a quite significant change to the asymmetry values, in all the cases, a clear $\sin\theta_{c.m.}$ behaviour appears in the BT asymmetry. In particular, it is the contact and t-channel kaon exchange in f_1 and f_3 that control the sign of $f_1 f_3^*$, and produce the positive BT asymmetry near threshold [14]. Also, note that it is natural that structures deviating from $\sin\theta_{c.m.}$ may arise at higher energies, since other mechanisms could become important, and f_1 and f_3 will no longer be the leading terms. Such a feature can be seen through the asymmetries for different K^* exchange phases at $W = 2.5$ GeV, where the interference of the $\cos\theta_{c.m.}(f_1 f_4^* - f_2 f_3^*)$ term in Eq. (10) shows up indeed.

Similar analysis can be applied to the production of $1/2^-$. In general, the transition amplitudes can be arranged in a way similar to the CGLN amplitudes for $1/2^+$:

$$\begin{aligned}\mathbf{J}\cdot\boldsymbol{\epsilon}_\gamma &= iC_1\frac{1}{|\mathbf{k}|}\boldsymbol{\sigma}\cdot(\boldsymbol{\epsilon}_\gamma\times\mathbf{k}) + C_2\frac{1}{|\mathbf{q}|}\boldsymbol{\sigma}\cdot\mathbf{q}\boldsymbol{\sigma}\cdot\boldsymbol{\epsilon}_\gamma \\ &\quad + iC_3\frac{1}{|\mathbf{q}||\mathbf{k}|^2}\boldsymbol{\sigma}\cdot\mathbf{k}\mathbf{q}\cdot(\boldsymbol{\epsilon}_\gamma\times\mathbf{k}) + iC_4\frac{1}{|\mathbf{q}|^2|\mathbf{k}|}\boldsymbol{\sigma}\cdot\mathbf{q}\mathbf{q}\cdot(\boldsymbol{\epsilon}_\gamma\times\mathbf{k}),\end{aligned} \tag{13}$$

where coefficients $C_{1,2,3,4}$ are functions of energies, momenta and scattering angle, and contain dynamical information on the transitions. Restricted to the kinematics near threshold, a term proportional to $\mathbf{q}\cdot\boldsymbol{\epsilon}_\gamma\boldsymbol{\sigma}\cdot(\mathbf{k}\times\mathbf{q})$ in the u-channel is neglected in the above expression, but included in the calculation. Such a term is the same order of C_4. However, they are both relatively suppressed in comparison with other terms due to the small $|\mathbf{q}|$ near threshold. As follows, we neglect the term of $\mathbf{q}\cdot\boldsymbol{\epsilon}_\gamma\boldsymbol{\sigma}\cdot(\mathbf{k}\times\mathbf{q})$ and express the above in parallel to the CGLN amplitudes.

Since $(\boldsymbol{\epsilon}_\gamma\times\hat{\mathbf{k}}) = i\lambda_\gamma\boldsymbol{\epsilon}_\gamma$, one can replace vector $(\boldsymbol{\epsilon}_\gamma\times\hat{\mathbf{k}})$ with $i\lambda_\gamma\boldsymbol{\epsilon}_\gamma$, and rewite the operator as

$$\begin{aligned}\mathbf{J}\cdot\boldsymbol{\epsilon}_\gamma = i\lambda_\gamma\Big[&iC_1\boldsymbol{\sigma}\cdot\boldsymbol{\epsilon}_\gamma + C_2\frac{1}{|\mathbf{q}||\mathbf{k}|}\boldsymbol{\sigma}\cdot\mathbf{q}\boldsymbol{\sigma}\cdot(\mathbf{k}\times\boldsymbol{\epsilon}_\gamma) \\ &+iC_3\frac{1}{|\mathbf{q}||\mathbf{k}|}\boldsymbol{\sigma}\cdot\mathbf{k}\mathbf{q}\cdot\boldsymbol{\epsilon}_\gamma + iC_4\frac{1}{|\mathbf{q}|^2}\boldsymbol{\sigma}\cdot\mathbf{q}\mathbf{q}\cdot\boldsymbol{\epsilon}_\gamma\Big],\end{aligned} \tag{14}$$

which has exactly the same form as Eq. (9) apart from an overall phase factor from the photon polarization $i\lambda_\gamma$. It also leads to the same form of the BT asymmetry as that for $1/2^+$:

$$D_{xz} = \sin\theta_{c.m.}\mathrm{Re}\{C_1C_3^* - C_2C_4^* + \cos\theta_{c.m.}(C_1C_4^* - C_2C_3^*)\}. \tag{15}$$

Quite remarkably, the behaviour of D_{xz} due to these two different parities now becomes more transparent since the role played by the dynamics has been isolated out. Nevertheless, it also results in vanishing asymmetries at $\theta_{c.m.} = 0°$ and $180°$.

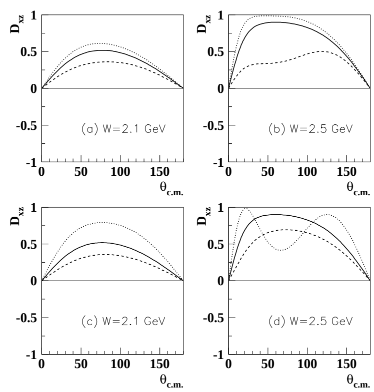

Figure 1. BT asymmetry for Θ^+ of $1/2^+$ at $W = 2.1$ and 2.5 GeV. The solid curves are results in the Born limit, while the dashed and dotted curves denote results with the K^* exchange included with different phases: $(g_{\Theta N K^*}, \kappa_\theta^*) = (-2.8, -3.71)$ (dashed curves in (a) and (b)), $(+2.8, +3.71)$ (dotted curves in (a) and (b)), $(-2.8, +3.71)$ (dashed curves in (c) and (d)), and $(+2.8, -3.71)$ (dotted curves in (c) and (d)).

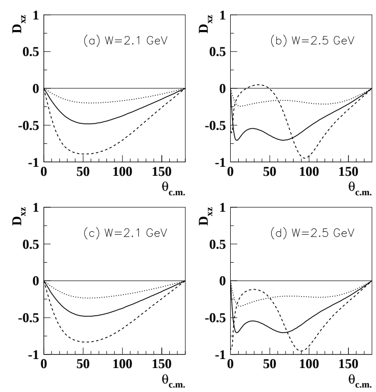

Figure 2. BT asymmetry for Θ^+ of $1/2^-$ at $W = 2.1$ and 2.5 GeV. The solid curves are results in the Born limit, while the dashed and dotted curves denote results with the K^* exchange included with different phases: $(g_{\Theta N K^*}, \kappa_\theta^*) = (-0.61, -0.371)$ (dashed curves in (a) and (b)), $(+0.61, +0.371)$ (dotted curves in (a) and (b)), $(-0.61, +0.371)$ (dashed curves in (c) and (d)), and $(+0.61, -0.371)$ (dotted curves in (c) and (d)).

The most important difference between the $1/2^-$ and $1/2^+$ cases is that the role of C_1 may not be as significant as f_1 in the production of $1/2^+$ due to the absence of the contact term in the production of $1/2^-$ state. In the Born limit, the main contribution to C_1, though dominant, comes from the s- and u-channel, which differs from the Kroll-Ruderman contribution to f_1 in $1/2^+$ production.

In Fig. 2, the BT asymmetries are calculated at $W = 2.1$ and 2.5 GeV in the Born limit and including the K^* exchanges. Interestingly, a rough $\sin\theta_{c.m.}$ behavior as the $1/2^+$ production still appears near threshold, but with different sign. A detailed analysis shows that near threshold, the term

of $C_1C_3^*$ is still the dominant one in the BT asymmetry. Here, the dominant contribution to C_1 is from the s-channel, while the dominant contribution to C_3^* is from the t-channel kaon exchange via the decomposition $\mathbf{q}\cdot\boldsymbol{\epsilon}_\gamma = \boldsymbol{\sigma}\cdot\mathbf{q}\boldsymbol{\sigma}\cdot\boldsymbol{\epsilon}_\gamma + i\boldsymbol{\sigma}\cdot(\boldsymbol{\epsilon}_\gamma\times\hat{\mathbf{k}})\mathbf{q}\cdot\hat{\mathbf{k}} - i\boldsymbol{\sigma}\cdot\hat{\mathbf{k}}\mathbf{q}\cdot(\boldsymbol{\epsilon}_\gamma\times\hat{\mathbf{k}})$. Therefore, we have

$$C_1C_3^* \simeq -e_0^2 g_{\Theta NK}^2 \frac{\kappa_n}{2M_n}\frac{2|\mathbf{k}||\mathbf{q}|}{(W-M_n)(t-M_K^2)}F_t(k,q)F_s(k,q), \qquad (16)$$

which will be negative since $\kappa_n = -1.91$. In comparison with Eq. (11), it gives a dynamical reason for the sign difference between these two parities. Meanwhile, the dominance of $C_1C_3^*$ only holds near threshold. With the increasing energy, other terms, e.g. $C_2C_4^*$, and other mechanisms, such as K^* exchange, can easily compete against $C_1C_3^*$ and produce deviations from the $\sin\theta_{c.m.}$ behavior as shown by the asymmetries at $W = 2.5$ GeV.

4. Discussions and summaries

In summary, we have analyzed the double polarization asymmetry, D_{xz}, in $\gamma n \to \Theta^+ K^-$, and showed it to be a useful filter for determining the parity of Θ^+, provided its spin-parity is either $1/2^+$ or $1/2^-$. Due to dynamical reasons, asymmetry D_{xz} near threshold would exhibit a similar behaviour but opposite sign. The advantage of studying polarization observables is that uncertainties arising from the unknown form factors can be partially avoided in a phenomenology. Therefore, although better knowledge of the form factors will improve the quantitative predictions, it should not change the threshold behaviour of D_{xz} dramatically. However, special caution should be given to the roles played by a possible spin-3/2 partner in the u-channel, as well as s-channel nucleon resonances. In particular, as studied by Dudek and Close [25], the spin-3/2 partner may have a mass close to the Θ^+. Thus, a significant contribution from the spin-3/2 pentaquark state may be possible. Its impact on the BT asymmetry needs to be investigated. In brief, due to the lack of knowledge in this area, any results for the BT asymmetry would be extremely important for progress in gaining insights into the nature of pentaquark states and dynamics for their productions [26]. Experimental facilities at Spring-8, JLab, ELSA, and ESRF should have access to the BT asymmetry observable.

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